

On the role of stochastic Fermi acceleration in setting the dissipation scale of turbulence in the interstellar medium

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ABSTRACT

We consider the dissipation by Fermi acceleration of magnetosonic turbulence in the Reynolds Layer of the interstellar medium. The scale in the cascade at which electron acceleration via stochastic Fermi acceleration (STFA) becomes comparable to further cascade of the turbulence defines the inner scale. For any magnetic turbulent spectra equal to or shallower than Goldreich-Sridhar this turns out to be $\geq 10^{12}$ cm, which is much larger than the shortest length scales observed in radio scintillation measurements. While STFA for such spectra then contradict models of scintillation which appeal directly to an extended, continuous turbulent cascade, such a separation of scales is consistent with the recent work of Boldyrev & Gwinn (2005) and Boldyrev & Konigl (2005) suggesting that interstellar scintillation may result from the passage of radio waves through the galactic distribution of thin ionized boundary surfaces of HII regions, rather than density variations from cascading turbulence. The presence of STFA dissipation also provides a mechanism for the non-ionizing heat source observed in the Reynolds Layer of the interstellar medium (Reynolds et al. 1999). STFA accommodates the proper heating power, and the input energy is rapidly thermalized within the low density Reynolds layer plasma.

1. Introduction

Radio scintillation has long been associated with interstellar turbulence (e.g. Minter and Spangler 1997)). A major requirement of turbulence-based scintillation models is that the inner scale of the cascade is comparable to the smallest scale of the scintillation. An alternative model has been proposed (Boldyrev & Gwinn 2003, 2005; Boldyrev & Konigl 2005) in which the scintillation instead results from index of refraction variations at the photoionized surfaces of a non-Gaussian distribution of either warm ISM regions or HII clouds. We find that Stochastic Fermi Acceleration (STFA) is a sufficiently efficient process

in the Reynolds Layer that it imposes a cutoff at a scale too large to be consistent with radio scintillation, supporting the non-Gaussian cloud model.

Observations of line ratios by Reynolds et al. (1999) using the Wisconsin H-alpha Mapper (WHAM) indicate that the Reynolds layer of the Milky Way interstellar medium (ISM) cannot be heated solely by photoionization. Variation with galactic latitude, $|z|$, of $[SII]/H\alpha$ and $[NII]/H\alpha$ both are found to be consistent with increases in electron temperature, T_e , but not in the ionization fractions of either species. Additionally, studies of external galaxies NGC891, NGC 4631, and NGC 3079 (Reynolds et.al. 2000; Otte et.al. 2001) show similar evidence of supplemental, non-photoionizing heat sources in the Reynolds layer, the region of low density, strongly ionized gas located ~ 1 kpc away from the galactic midplane. These studies supplement the $[SII]/H\alpha$ and $[NII]/H\alpha$ ratios with analysis of $[OIII]/H\beta$ and $[OII]/H\alpha$, which also are consistent with non-ionizing heating. The heating rate in the Milky Way is consistent with being either proportional to electron density, n_e , or density independent. The heat source is empirically found to have a power input of $\epsilon = G_1 n_e$, where $G_1 \sim 10^{-25}$ erg s $^{-1}$, or $\epsilon = G_2$, where $G_2 \sim 10^{-27}$ erg s $^{-1}$ cm $^{-3}$. This allows the supplemental heat source to dominate at high $|z|$, where n_e is low and photoionization heating to dominate at low $|z|$ where n_e is high. A number of potential mechanisms are presented by Reynolds et al. (1999) as supplemental heat sources, including photoelectric grain heating, Coulomb collisions with cosmic rays, magnetic field reconnection, and dissipation of superbubble driven magnetic turbulence.

We consider the role of STFA in the dissipation of interstellar turbulence, looking primarily for the inner scale of the cascade. ISM turbulence is driven principally by superbubbles, the large blast shells carved out when the members of an OB association reach the supernova stage in a short time period. STFA occurs when electrons traveling in along a magnetic field line encounter moving compressions of the local magnetic field and are reflected. While STFA drains turbulence primarily into ions for thermal pressure dominated plasmas, electrons are the primary energy recipient when the ion speed is sub-Alfvénic (i.e. LaRosa, et al. (1996); Blackman (1999); Selkowitz & Blackman (2004)). The latter case is relevant for the present study because the Reynolds layer seems to be a weakly magnetically dominated plasma (Minter & Spangler 1996).

The energy change, δE from a single STFA mirroring is (Fermi 1949; Achterberg 1981; Longair 1994; Selkowitz & Blackman 2004)

$$\frac{2E}{c^2}(v_A^2 \pm v_e v_A), \quad (1)$$

where E is the initial energy of the reflected particle, c is the speed of light, v_A is the Alfvén

speed, and v_e is the electron speed. Typical Reynolds layer temperatures are $0.6 < T_4 < 1.2$ where $T_4 = T_e/10^4\text{K}$. This corresponds to an energy in the range $0.5 < E < 1.1\text{eV}$. Thus δE is significantly below ionization energies, which are tens of eV, and even if the reflection rate is high, STFA of protons is a non-ionizing process in the turbulent Reynolds layer of the ISM.

In section 2 we show that STFA provides the correct power and that the energy input is quickly thermalized, consistent with the conditions imposed on the heating source. We then determine the truncation scale for the turbulent cascade. In section 3, we discuss the implications of the inferred STFA dissipation scale for models of interstellar scintillation, finding that STFA truncates the cascade at a long length scale, inconsistent with turbulence models of radio scintillation. We conclude in section 4.

2. Stochastic acceleration in the turbulent Reynolds layer

The power available for electron heating in the MHD turbulent cascade, ϵ_T , can be estimated as

$$\epsilon_T = n_e m_p \frac{v_A^3}{L}, \quad (2)$$

where n_e is the electron number density, v_A is the local Alfvén speed, and L is the outer scale of the interstellar turbulence inertial range. From Minter and Spangler (1997): $v_A = 2.3 \times 10^6 \text{ cm s}^{-1}$ and $L = 10^{19} \text{ cm}$. Minter & Spangler (1996) infer this value for L by analysis of emission measure and rotation measure structure functions. From Reynolds et al. (1999): $n_e = 0.28 \text{ cm}^{-3}$ at galactic latitude $|z| = 1 \text{ kpc}$. $\epsilon_T = 5 \times 10^{-26} \text{ erg s}^{-1} \text{ cm}^{-3}$, consistent with the heating rate called for in Reynolds et al. (1999) and calculations of superbubble injected power (Mac Low, & Klessen 2004; Elmegreen & Klessen 2004).

Selkowitz & Blackman (2004) examined STFA in the $v_A \gg c_s$ regime with applications to solar flares, and concluded that the post-acceleration spectrum is dominated by the rate of escape from the acceleration region. The post-escape spectrum is important in flares because the observed x-ray emission is generated via Bremsstrahlung as electrons escape the acceleration region and encounter the dense chromospheric plasma. In the turbulent Reynolds layer, the acceleration region is effectively infinite in extent. Instead of observing emission by escaped electrons, the observed emission comes from electrons which remain in the Reynolds layer. In this case we are concerned with the spectrum of electrons still confined within the acceleration region. Regardless, the acceleration rate found in Selkowitz & Blackman (2004) remains valid

$$\left(\frac{dE}{dt}\right)_s = \langle \delta E \rangle R, \quad (3)$$

where

$$\langle \delta E \rangle = 4m_e v_A^2, \quad (4)$$

is the average energy per reflection and

$$R = \frac{v_d}{2L} \left(\frac{\lambda_s}{L}\right)^2, \quad (5)$$

is the rate of reflections, such that

$$\left(\frac{dE}{dt}\right)_s = \frac{2}{L} m_e v_A^2 v_d \left(\frac{\lambda_s}{L}\right)^{-(1-1/a)}, \quad (6)$$

and λ_s is the turbulent inner scale (associated with the parallel component of the magnetic field gradient for anisotropic turbulence), v_d is the effective relative speed of the electrons and compressions within the plasma, m_e is the electron mass, and a is the spectral index of the turbulent cascade

$$\frac{\lambda}{L} = \left(\frac{\delta B}{B}\right)^a. \quad (7)$$

It is assumed throughout our analysis that at length scales comparable to L the strength of the magnetic fluctuations, δB is comparable to that of the mean magnetic field, B . The magnetic spectrum of turbulence in the ISM is difficult to measure. Between 10^{19} pc and 0.03pc the spectrum appears to be close to Kolmogorov ($a = 3$) (and flatter on larger scales) (Han et al. (2004); Minter & Spangler (1996)). There is no direct measure of the magnetic spectrum on scales below 0.03pc.

For the range of magnetic spectra $2 \leq a \leq \infty$ the result from our calculations to follow, that the cascade truncation scale well exceeds the smallest scintillation scale, will not change. However, we first consider a Goldreich-Sridhar (hereafter GS) turbulent spectrum with $a = 2$ (Goldreich and Sridhar 1997) (see also Matthaeus et al. (1998)). Like Matthaeus et al. (1998), GS turbulence is fully anisotropic, and is based on an incompressible cascade. It involves a more rapid cascade in the direction perpendicular to the local mean field than parallel to it. Turbulence becomes less and less compressible on smaller scales, and since the Reynolds layer seems to be modestly magnetically dominated (Minter & Spangler (1996)), GS is plausible on small enough scales: The GS spectrum more closely arises in compressible and incompressible simulations when an initially relative strong field is imposed (e.g. Maron & Goldreich (2001), Cho et al. (2002a), Vestuto et al. (2003), Haugen & Brandenburg (2004b)), i.e. when the initial ratio of thermal to magnetic pressure, $\beta < 1$ and the velocity

fluctuations are of order the initial field. Other values for a should not be ruled out however, because the magnetic turbulent spectrum for galactic ISM conditions is not a solved theoretical problem. In particular, when an initially weak field is imposed, driven incompressible simulations show that the magnetic spectral index may be flatter, and flatter than the velocity spectra. Although it may approach 5/3 for magnetic Prandtl number of order unity (Haugen et al. (2003)) there may be a trend toward further flattening at larger magnetic Prandtl numbers (Schekochihin et al. (2002); Maron et al. (2004); Schekochihin et al. (2004); Haugen et al. (2004a)). The magnetic spectrum is steeper in the presence of kinetic helicity (Maron & Blackman (2002)). The compressible driven simulations of Vestuto et al. (2003) show that the stronger the initial field, the closer the magnetic and velocity spectra match, and the flatter the magnetic spectra the weaker the initial field. The role of the magnetic Prandtl number is hard to assess in Vestuto et al. (2003) since no explicit viscosity or resistivity is used.

Proceeding with $a = 2$, we note that for this case, Selkowitz & Blackman (2004) find that the parallel cascade law is relevant to STFA, and λ in eq. (7) corresponds to the parallel cascade law. When electrons are able to freely stream from one STFA scattering site to the next without deflection by pitch angle scattering (the free streaming limit) v_d is equal to the electron speed v_e and,

$$\left(\frac{dE}{dt}\right)_f = \frac{2}{L} m_e v_A^2 v_e \left(\frac{\lambda_f}{L}\right)^{-1/2}. \quad (8)$$

In order to determine the truncation scale of the cascade, we impose the balance condition that the turbulent power, ϵ_T is equal to the STFA acceleration rate. Setting the STFA rate equal to the total power selects a single value of λ_f , which is an effective upper limit on the truncation scale. We use the subscript f to denote the free streaming limit.

In the limit of very strong pitch angle scattering, where the scattering length scale $\lambda_p \ll \lambda_f$, λ_p can be considered the electron mean free path. Electrons are effectively trapped in regions smaller than λ_f , executing a random walk with small drift speed. They encounter magnetic compressions only as the compressions stream past, significantly reducing the acceleration rate below that of Eq 8. The power balance condition between STFA and the cascade power is not met, and the cascade will proceed to scales shorter than λ_f . The cascade cannot continue farther than a length scale comparable to λ_p . At this point, the system passes back into a free-streaming regime, with a substantially shorter λ than is required to meet the power balance condition. The turbulence then drains rapidly.

From the above two paragraphs, we conclude that the cascade would be truncated at a scale no smaller than the minimum of λ_p and λ_f . To assess the minimum scale of dissipation

for the Reynolds layer, we assume that the pitch angle scattering is dominated by electron-electron Coulomb scattering. Spitzer (1967) finds the Coulomb self-collision time, t_c , to be

$$t_c = \frac{0.266}{n_e} \frac{T_e^{3/2}}{\ln \Lambda}, \quad (9)$$

where T_e is the electron temperature, and $\ln \Lambda \sim 25$ is determined by the effective long range cutoff of the Coulomb force in a plasma. The electron speed is given by $v_e = \sqrt{2kT/m_e}$, and thus

$$\lambda_p = v_e t_c = \sqrt{\frac{2kT}{m_e}} \frac{0.266}{n_e} \frac{T_e^{3/2}}{\ln \Lambda} = 2 \times 10^{13} T_4^2, \quad (10)$$

where $T_4 = T_e/10^4$. For the range of observed temperatures in Reynolds et al. (1999), $0.6 < T_4 < 1.2$, the electron mean free path falls in the range $8 \times 10^{12} \text{cm} < \lambda_p < 3 \times 10^{13} \text{cm}$. Taking the lower limit of this range fixes the lower bound of the cascade truncation scale.

In the free streaming limit we can find λ_f by setting $\epsilon_T = n_e (\frac{dE}{dt})_f$,

$$\lambda_f = \left(\frac{2m_e}{m_p} \right)^2 \left(\frac{v_e}{v_A} \right)^2 L = 6 \times 10^{14} T_4^{-1}. \quad (11)$$

For the observed temperature range of $0.6 < T_4 < 1.2$, $10^{15} \text{cm} > \lambda_f > 5 \times 10^{14} \text{cm}$. This sets an upper bound on the cascade truncation scale.

It should be noted that the free streaming approximation is invalid in the Reynolds layer as $\lambda_f/\lambda_p = 100$ at $T_4 = 6$. There are, on average 100 pitch angle scattering events per electron per encounter with a compression. Given the low fraction of encounters which result in a reflection, F , electrons cannot stream freely from one reflection site to the next. The quantity F is determined by the pitch angle condition for reflection

$$\cos^2 \theta_{min} < \frac{\delta B}{B}, \quad (12)$$

and the turbulent cascade law (e.g. eq. (7) for GS turbulence). All electrons with $\theta > \theta_{min}$ reflect. Assuming pitch angle isotropy, we estimate

$$F = \frac{1}{4\pi} \int_{\theta_{min}}^{\pi} 2\pi \sin(\theta) d\theta = \frac{1}{2} (\lambda_f/L)^{1/4}, \quad (13)$$

where we have integrated over all angles greater than θ_{min} to find the fraction of phase space which satisfies the pitch angle condition for reflection, and we have used (7) for the last

equality. Electrons which have too large a component of their momentum in the direction parallel to the field are not stopped by the compressions, and pass through them. The acceleration rate must be retarded significantly. We have bounded the turbulent dissipation scale $8 \times 10^{12} \text{cm} < \lambda_s < 10^{15} \text{cm}$ for GS turbulence for a Goldreich-Sridhar, $a = 2$, cascade.

In order to produce a dissipation scale consistent with the smallest scales observed in scintillation measurements, the cascade must truncate at the inner scale predicted by scintillation models. Stated values of this scale vary: $3.5 \times 10^6 \text{cm}$ in Moran, et.al. (1990), $3 \times 10^7 \text{cm}$ in Molnar, et.al. (1995), and $3 \times 10^{10} \text{cm}$ in Rickett (1990). Rewriting eq 11 for an arbitrary value of a , we have

$$\lambda_f = \left(\frac{2m_e}{m_p} \right)^{1-1/a} \left(\frac{v_e}{v_A} \right)^{1-1/a}. \quad (14)$$

In order to produce a value $\lambda_f = 10^{10} \text{cm}$, the cascade must have a very steep spectrum, with $a < 1.3$. The other predictions require even steeper spectra. Such spectra are inconsistent with simulations of interstellar turbulence, as well as the inferred spectral indices of scintillation models, which place the lower limit near a Kolmogorov spectrum, $a = 3$ (Han et al. 2004; Minter and Spangler 1997).

Assuming the Reynolds layer is weakly magnetically dominated (Minter & Spangler 1996) implies that we have studied dissipation of fast magnetosonic mode turbulence. On the other hand, Lithwick & Goldreich (2001) argue that for thermally dominated plasmas, $\beta > 1$, the turbulent density fluctuations are dominated by the slow and entropy mode, and that the fast mode is decoupled from the slow, entropy, and Alfvén modes. Of the slow and entropy modes, only the slow mode is important for STFA. While the fast mode could be dominant for STFA in the Reynolds layer, the acceleration by slow modes does achieve the same qualitative result, albeit at a shorter dissipation scale. In the Reynolds layer, where $\beta \sim 0.1$, the slow mode velocity is given roughly by the sound speed $c_s \sim 0.2v_A$. If both modes are present with equal energy density, the fast mode is more efficient for STFA than the slow mode. To apply eq. (6) to STFA of slow modes in a plasma with $\beta < 1$ requires the replacement of v_A with c_s , such that

$$\left(\frac{dE}{dt} \right)_s = \frac{2}{L} m_e c_s^2 v_d \left(\frac{\lambda_s}{L} \right)^{-(1-1/a)}. \quad (15)$$

For free streaming electrons, there is no change in $v_d = v_e$ and the dissipation scale is only decreased by a factor of $(c_s/v_A)^2 = 0.04$. For trapped electrons, the mean free path continues to place a lower bound on the dissipation. Despite the moderate decrease in λ_s in a slow mode dominated cascade, it remains well above the scintillation length scale.

In addition to providing the correct power input and a physically plausible dissipative scale, STFA must supply energy in a manner consistent with heating, as opposed to non-thermal acceleration. Selkowitz & Blackman (2004) demonstrated that, in the limit of no self-interaction which is appropriate for solar flare applications above a few keV, the electron energy spectrum within the acceleration region is shaped by two competing effects: a bulk shifting to higher energy, and a diffusive spreading of the spectrum. The resulting spectrum can be quasi-thermal, even when electrons are not truly sharing energy. However, in the case of ISM heating, where Coulomb self-scattering of electrons is presumed to be the dominant source of pitch angle scattering, STFA is accompanied by rapid thermalization of the electron population. Consider the acceleration time, τ_{STFA}

$$\tau_{STFA} = \frac{E}{\left(\frac{dE}{dt}\right)_s} = 3 \times 10^{12} T_4 s, \quad (16)$$

and the thermalization time, τ_{eq} , (Spitzer 1967)

$$\tau_{eq} = 6 \times 10^5 T_4^{3/2} s, \quad (17)$$

where we have taken the thermalization time to be the two species equilibration time, with both species being electrons at a single temperature, T_4 . When $\tau_{eq} < \tau_{STFA}$, as is the case for the Reynolds layer, then the energy input via STFA is shared among electrons rapidly; the energy spectrum is thermal. The electron-proton thermalization time is greater than τ_{eq} by a factor of the mass proton to electron ratio $m_p/m_e = 1836$, which is nevertheless shorter than the acceleration time. The protons (and heavy ions) remain in equilibrium with the electrons as well.

3. Implications of the turbulent cutoff scale for models of interstellar scintillation

The dominant power source for interstellar turbulence is likely supernovae and superbubble shells (Elmegreen & Klessen (2004); Mac Low, & Klessen (2004)). Superbubbles are powerful and frequent enough to periodically pass through, and thus reseed turbulence in, the entire galactic disk. If the efficiency for converting a single supernova's mechanical luminosity of 10^{51} erg to turbulent energy is $\eta_{SN} = 0.1$, the rate of supernova events in the Milky Way is $R_{SN} = 50 \text{ yr}^{-1}$, then the total turbulent driving from supernovae and superbubbles is estimated to be $3 \times 10^{-26} \text{ erg s}^{-1}$ (reviews by (Mac Low, & Klessen 2004; Elmegreen & Klessen 2004) and references within). This is consistent with the estimated damping rate ϵ_T from eq (1) as well as the required supplemental heating rate (Reynolds et al. 1999).

Although we have considered a GS cascade law because it is inherently anisotropic, and incorporates the magnetic field into the formalism directly, we emphasize again that the GS cascade, with $a = 2$ magnetic spectral index lies at one end of a range of possible cascade power laws (e.g. Goldreich and Sridhar (1997); Vestuto et al. (2003)). As discussed, varying a over a wide range does not substantially alter the results of our analysis.

Likewise, it is not certain which magnetosonic wave mode, fast or slow, is dominant in the Reynolds layer. Lithwick & Goldreich (2001) argues for slow modes to dominate in a high β plasma. Fast mode turbulence may play a stronger role in $\beta \sim 0.1$ diffuse Reynolds layer plasma. Because the fast mode is a more efficient STFA accelerator, due primarily to its higher phase speed, we emphasized STFA on fast modes, but it was shown in section 2 that the major qualitative result of our study holds for either wave mode: STFA is a sufficient source of heating, and the dissipation scale of MHD turbulence is longer than the short scales observed in radio scintillation measurements.

The correspondence between the turbulent driving power and the inferred supplemental heating rate of Reynolds et al. (1999) is suggestive of a connection. This connection would be more strongly verifiable if the turbulent dissipation scale could be linked to direct observations of the Reynolds layer. This has not yet been done. Spangler (1991) and Minter and Spangler (1997) considered that the radio scintillation observations (e.g. Armstrong, Rickett, Spangler 1995) may be consistent with scattering by localized turbulent density fluctuations across the inertial range of Kolmogorov turbulence, but they further suggest that the “fluctiferous” medium is either HII region envelopes, or the higher density portions of the warm ionized medium (WIM). Although possible turbulent dissipation scales (like the Larmor radius) of the scintillating region correspond well with the smallest scales of radio scintillation, no scintillation measurements are analyzed in lines of sight through the low density Reynolds layer. Minter and Spangler (1997) instead assume that either the Larmor radius or the ion inertial length is the key dissipative scale in the Reynolds layer, just as in denser regions of the ISM, and use this assumption to calculate dissipation and heating rates.

We find that the STFA turbulent dissipation scale is instead likely far larger ($\lambda_s \geq 10^{13}\text{cm}$) in the Reynolds layer, for GS turbulence. The STFA dissipation scale does drop to the observed scintillation scale in the case of a very steep turbulent spectrum at small scales, with index $a \leq 1.3$. Interpretations of the scintillation as turbulence driven, such as Minter and Spangler (1997), indicate a spectrum no steeper than Kolmogorov, $a = 3$. Although STFA is seemingly at odds with radio scintillation measurements, it may not be if the source of the scintillation is not a continuous turbulent cascade as is commonly assumed. In particular, recent studies have found that interstellar radio scintillation is in fact consistent

with a non-Gaussian distribution of stellar ionization shells Boldyrev & Gwinn (2003, 2005); Boldyrev & Konigl (2005). Local index of refraction variations at the thin photoionized surfaces of molecular clouds are distributed randomly throughout the interstellar medium. Radio waves are deflected through small angles at these fronts, and adjacent rays can be deflected through differing paths. A detailed model employing Levy statistics provides an excellent fit to the temporal spreading of pulses from distant pulsars, which traditional cascade scintillation models have had difficulty explaining. Boldyrev & Gwinn (2005) and Boldyrev & Konigl (2005) conclude that the density fluctuations are more likely to result from randomly distributed thin ionization shells in the ISM than from global turbulence. As a result, the scintillation would no longer provide a constraint on the inner scale of the turbulent cascade. Turbulence driven on very large (≥ 10 pc) scales could damp on scales larger than the smallest scintillation scales and dissipation of the turbulence by STFA would not conflict with scintillation measurements.

The lower bound of the dissipation scale found above violates an assumption used in the approach of Minter and Spangler (1997). Given that the measured supplemental heating rate in the low density ISM can be well explained by turbulent dissipation, one may infer that turbulence is driven and dissipated throughout the ISM. The density is higher and temperature lower in the lower $|z|$, and thus λ_p is shorter, in the portions of the WIM where the scintillation may occur. STFA is then even more closely tied to the lower bound for λ_s set by Coulomb scattering at low $|z|$. If STFA is important, it would then be unlikely that turbulence input by superbubbles cascades to sufficiently small scales in any region of the interstellar medium to account for the scintillation observations. This gives credence to the argument of Boldyrev & Gwinn (2003, 2005); Boldyrev & Konigl (2005), which does not rely on cascading turbulence but concludes that the statistical spatial distributions of the ionized boundaries of molecular clouds better explains scintillation than do models based on turbulence dominated density fluctuations.

4. Conclusion

Dissipation of turbulence almost certainly plays a role in the heating of the interstellar medium. The turbulent energy supply from supernovae is sufficient to provide the supplemental heating source required by the observations of Reynolds et al. (1999); Reynolds et.al. (2000); Otte et.al. (2001). We have shown that STFA can act as the damping mechanism in the Reynolds layer of the Milky Way, truncating the turbulent cascade at length scales no shorter than 8×10^{13} cm for a cascade no steeper than Goldreich-Sridhar. This truncation scale is too large to be consistent with turbulence-based models of interstellar scintillation.

Instead, models like those of Boldyrev & Gwinn (2003, 2005); Boldyrev & Konigl (2005), which do not rely on the cascade, are required.

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